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SOME PHYSICS OF THE GRUNEISEN PARAMETER

Paul Harris, et al

Picatinny Arsenal  
Dover, New Jersey

September 1972

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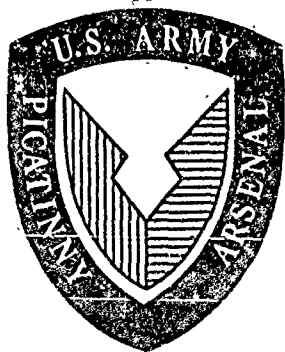
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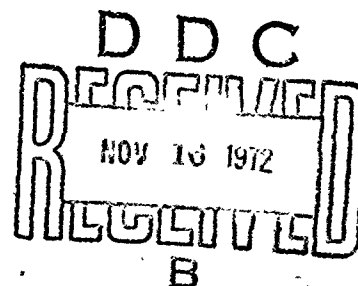
PAUL HARRIS  
LOUIS AVRAMI

SEPTEMBER 1972

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## DOCUMENT CONTROL DATA - R &amp; D

(Security classification of title, body of abstract and indexing annotation must be entered when the overall report is classified).

1. ORIGINATING ACTIVITY (Corporate author) U.S. Army, Picatinny Arsenal, Dover, N.J. 07801		2a. REPORT SECURITY CLASSIFICATION UNCLASSIFIED	
		2b. GROUP N/A	
3. REPORT TITLE  Some Physics of the Grüneisen Parameter			
4. DESCRIPTIVE NOTES (Type of report and inclusive dates) The physics of the Grüneisen parameter (thermomechanical coefficient) is examined.			
5. AUTHOR(S) (First name, middle initial, last name) Paul Harris and Louis Avrami			
6. REPORT DATE September 1972		7a. TOTAL NO. OF PAGES 39	7b. NO. OF REFS 18
8a. CONTRACT OR GRANT NO. N/A		9a. ORIGINATOR'S REPORT NUMBER(S) Technical Report 4423	
b. PROJECT NO.		9b. OTHER REPORT NO(S) (Any other numbers that may be assigned this report)	
c.		N/A	
d.			
10. DISTRIBUTION STATEMENT  Approved for public release; distribution unlimited.			
11. SUPPLEMENTARY NOTES		12. SPONSORING MILITARY ACTIVITY U.S. Army Picatinny Arsenal, Dover, N.J.	
13. ABSTRACT  The physics of the Grüneisen parameter for solids is investigated from the view of conceptual and physical contributions. The relatively simple physics of a van der Waal's gas is applied to solid densities yielding a Grüneisen $\Gamma \approx 2$ .  The separate electron and phonon contributions for metals are considered, and it is shown that except for very high and very low temperatures the phonon contribution dominates. Procedures for evaluating $\Gamma$ from shock-Hugoniot data for isotropic solids are treated, and the necessity for applying ultrasonic data to anisotropic crystals is considered.  The role of the more exotic thermal excitations of solid state physics in determining $\Gamma$ is mentioned.			

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OBSOLETE FOR ARMY USE.

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KEY WORDS

LINK A

LINK B

LINK C

ROLE

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*II*

## TABLE OF CONTENTS

	<u>PAGE NO.</u>
LIST OF TABLES AND FIGURES	1
ABSTRACT	2
SECTION	
I. INTRODUCTION	3
II. GASES	5
A. Perfect gases	
B. Imperfect gases	
III. GENERAL THEORY	10
A. Thermomechanical Considerations	
B. Statistical Mechanics Considerations	
C. General Applications	
IV. SOLIDS	20
A. Electrons	
B. Phonons - Low Temperatures	
C. Phonons - High Temperatures	
D. Volume Dependence	
E. Ultrasonic Data	
V. DISCUSSION AND SUMMARY	30
REFERENCES	33
DISTRIBUTION LIST	34

## TABLE OF CONTENTS

	<u>PAGE NO.</u>
LIST OF TABLES AND FIGURES	1
ABSTRACT	2
SECTION	
I. INTRODUCTION	3
II. GASES	5
A. Perfect gases	
B. Imperfect gases	
III. GENERAL THEORY	10
A. Thermomechanical Considerations	
B. Statistical Mechanics Considerations	
C. General Applications	
IV. SOLIDS	20
A. Electrons	
B. Phonons - Low Temperatures	
C. Phonons - High Temperatures	
D. Volume Dependence	
E. Ultrasonic Data	
V. DISCUSSION AND SUMMARY	30
REFERENCES	33
DISTRIBUTION LIST	34



## LIST OF TABLES

### TABLE NO.

### PAGE NO.

- I - Comparison of Shock and  
Thermodynamic Data

25

### FIGURE NO.

- I - Gruneisen Parameter as a function of time  
for beam pulse width time  $t$ . The time for  
translation between vibrational states is  $t_1$ .

6

# ABSTRACT

The physics of the Grüneisen parameter for solids is investigated from the view of conceptual and physical contributions. The relatively simple physics of a van der Waal's gas is applied to solid densities yielding a Grüneisen  $\Gamma \approx 2$ .

The separate electron and phonon contributions for metals are considered, and it is shown that except for very high and very low temperatures the phonon contribution dominates. Procedures for evaluating  $\Gamma$  from shock-Hugoniot data for isotropic solids are treated, and the necessity for applying ultrasonic data to anisotropic crystals is considered.

The role of the more exotic thermal excitations of solid state physics in determining  $\Gamma$  is mentioned.

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The physics of the Grüneisen parameter for solids is investigated from the view of conceptual and physical contributions. The relatively simple physics of a van der Waal's gas is applied to solid densities yielding a Grüneisen  $\Gamma \approx 2$ .

The separate electron and phonon contributions for metals are considered, and it is shown that except for very high and very low temperatures the phonon contribution dominates. Procedures for evaluating  $\Gamma$  from shock-Hugoniot data for isotropic solids are treated, and the necessity for applying ultrasonic data to anisotropic crystals is considered.

The role of the more exotic thermal excitations of solid state physics in determining  $\Gamma$  is mentioned.

## I. INTRODUCTION

Consider a system having total internal energy  $E$ , pressure  $P$ , volume  $V$ , and temperature  $T$ . The Grüneisen parameter,  $\Gamma$ , for that system is defined by

$$\Gamma \equiv V \left( \frac{\partial P}{\partial E} \right)_V, \quad (1)$$

where the derivative is to be evaluated at constant volume. As given by the macroscopic (operational) definition of Eq. (1),  $\Gamma$  is a measure of the change in pressure produced by a change in system total energy under the condition of constant volume.

As evidenced from the above definition  $\Gamma$  is a thermo (energy)-mechanical (pressure) quantity, and as such is expected to be important to thermomechanical problems. A few examples of such problems are: shock wave effects (e.g., explosively driven shock waves in solids), the thermal expansion of solids, and the rapid heating of materials due to their absorbing intense pulses of nuclear radiation. In the last example  $\Gamma$  becomes of extreme importance since, for radiation pulse duration small compared to relevant times for acoustic transport, the induced thermal pressure is directly proportional to  $\Gamma$ .

In this report we shall consider the physics, sometimes on a microscopic scale, responsible for the Grüneisen parameter having its observed values. There is good reason for doing so. Presently, in the United States, technology has outpaced science to the point where design engineers frequently find themselves using materials which have unknown responses to the environments in which they are to be placed. This is especially true in the area of military

missile technology where very often new exotic materials are required to survive radiation environments in which their response is unknown. Knowledge of the Grüneisen constant physics would allow an engineer to make valid estimates for those materials for which  $\Gamma$  has not yet been measured.

Understanding the physics of  $\Gamma$  is also quite important to the understanding of related physics problems. One such problem area of interest is that of electromechanical shock effects in metals. If one analyzes shock propagation in metals in terms of a two fluid model of electrons and phonons, then the Grüneisen parameter of each component becomes important. Some of the analysis presented in this report will be directed towards that problem.

So that the uninitiated reader may be eased gently into the problem we will begin by treating the relatively simple problem of the Grüneisen parameter of gases. After that we will consider some generalities of  $\Gamma$ , in particular as they apply to solids. Finally, calculations will be made of the Grüneisen parameter for some real systems of interest.

## II. GASES

### A. Perfect Gases.

Consider an ideal monatomic gas with equation of state (EOS) and total energy  $E$  given by

$$PV = NkT, \quad E = \frac{3}{2} NkT, \quad (2)$$

$$\therefore PV = \frac{2}{3} E. \quad (3)$$

From Eqs. (1) and (3) we find

$$\Gamma(\text{ideal gas}) = \frac{2}{3}. \quad (4)$$

In the above  $N$  is the number of molecules (atoms for a monatomic gas),  $k$  is the Boltzmann constant ( $1.38 \times 10^{-16}$  ergs/°Kelvin), and  $T$  is the absolute temperature.

Looking now at a gas made up of diatomic molecules with  $E = \frac{5}{2} NkT$  we find  $\Gamma = \frac{2}{5}$ . Consequently, a given amount of energy introduced into a diatomic gas produces less of a pressure increase than in the corresponding monatomic gas. This result is easy to understand; for the diatomic case some of the energy goes into the internal energy associated with molecular vibration.

While the above example is somewhat trivial there is information of interest contained in it because the diatomic gas has a hint of solid state behavior in it by virtue of its internal energy. Suppose that we were able to put the energy into the diatomic gas system in a time small compared to that necessary for transitions

between the translational states of a diatomic molecule. Since  $E = \frac{5}{2} NkT$  is composed of  $\frac{3}{2} NkT$  kinetic energy associated with translational motion of the center of mass of each molecule, and  $NkT$  internal energy associated with vibrational states, the  $\Gamma$  for the system would appear as shown in Fig. 1 below.

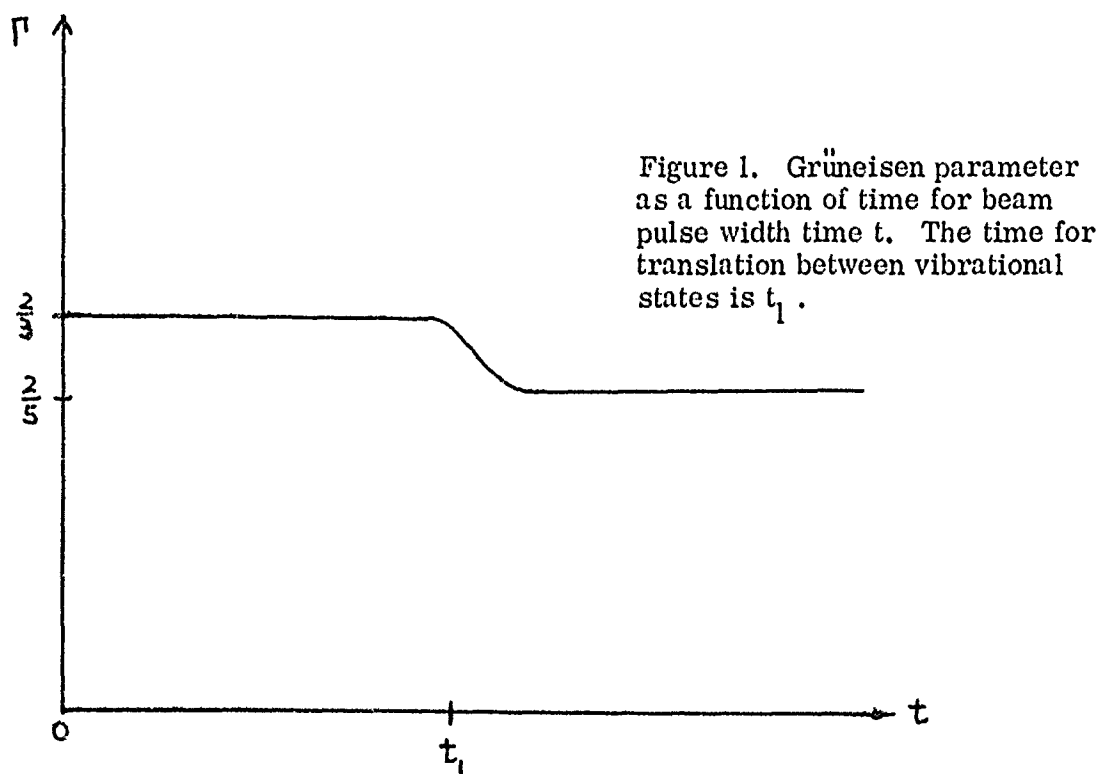


Fig. 1 shows the  $\Gamma$  that would be measured as a function of varying deposition time for the system of diatomic molecules. It is to be understood that we are only talking about a thought experiment here since  $t_1 \sim 10^{-14}$  sec. As an alternative to the fast energy pulse the same results (i.e.,  $\Gamma = \frac{2}{3}$  for a diatomic gas system) can be achieved by imagining an omnipotent hand reaching in and maintaining the pre-radiation separation distance, and thus the pre-radiation state, between the atoms of each molecule.

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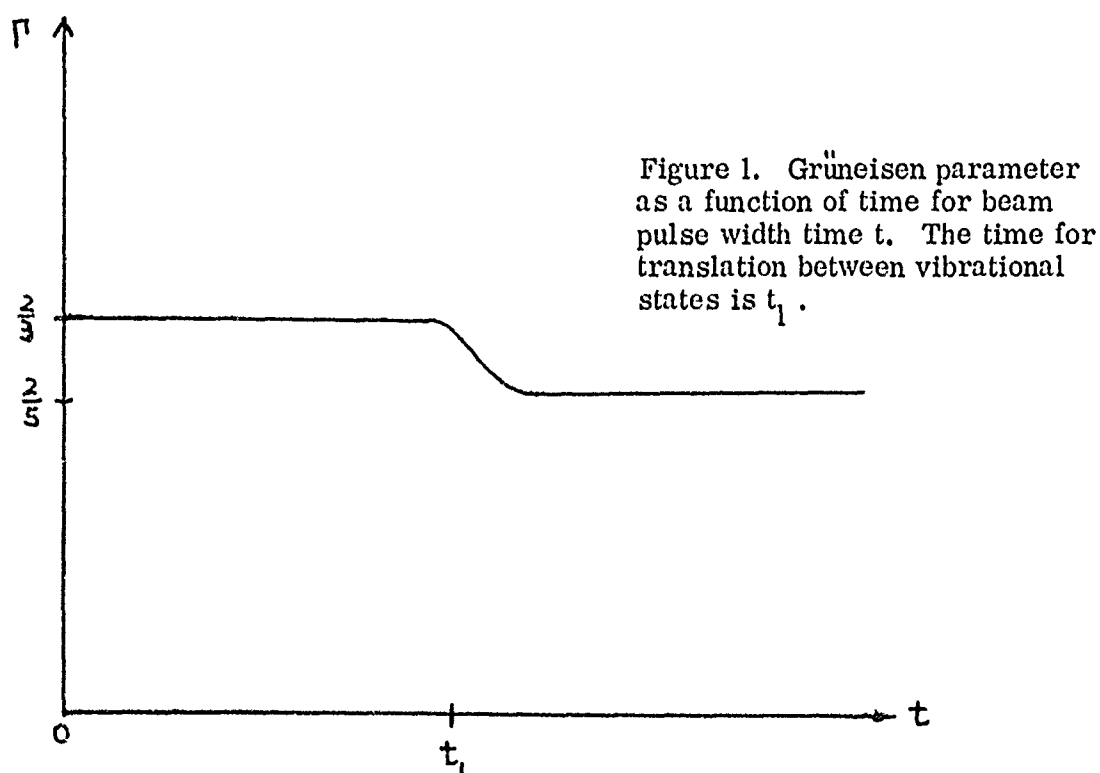


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(other than the approximately hard sphere interaction responsible for the volume exclusion constant  $b$ ) does not appear in  $\Gamma$ . This is simply understood; at constant volume the average intermolecular distance remains unchanged by the addition of energy, and thus the potential energy from intermolecular interactions remains unchanged. We are thus led to expect that the intermolecular potential will not appear in the constant volume Gruneisen constant of a solid. Such is indeed the case.

Suprisingly enough Eq. (8) can be used to estimate the  $\Gamma$  of a liquid or a solid. One simple model for  $b$  has  $b = 4V_0$ , where  $V_0$  is the volume of a molecule. Applying Eq. (8) to a solid or liquid with half of its available volume occupied by the molecules (and half otherwise empty), with the simple model for  $b$  mentioned above, yields

$$\Gamma = \frac{2}{3} \left( 1 + \frac{4}{2} \right) = 2, \quad (9)$$

which is a suprisingly acceptable result; with the exception of Zirconium, metals have Gruneisen constants<sup>2</sup> which vary between one and three at room temperature and atmospheric pressure.

It is possible to understand the success of Eq. (9), as applied to a solid or a liquid, by considering the shock compression Hugoniot<sup>3</sup> EOS for a condensed medium. For volume changes small enough so that only the first nonlinear term need be considered we can write

$$P_H = A_H \left( \frac{P}{P_0} - 1 \right) + B_H \left( \frac{P}{P_0} - 1 \right)^2, \quad (10a)$$

or upon rewriting

$$P_H V = A'_H V_0 + B'_H \frac{V_0^2}{V} + C', \quad (10b)$$

where the subscript zero denotes the undeformed state, the subscript H denotes Hugoniot, and  $A_H$ ,  $B_H$  are characteristic of the material in question.  $\rho$  is the mass density, and the primed quantities are the constants associated with the rearrangement in Eq. (10b). The point of interest is that Eqs. (5) and (10b) have the same functional form ( $A_H$  is a function of temperature), and thus the proper interpretation of  $b$  in Eq. (8) should lead to a proper result -- which it does.

### III. GENERAL THEORY

There are two ways in which to approach the problem of the Gruneisen parameter. One method involves purely macroscopic thermomechanical considerations, while the other involves a detailed statistical mechanics investigation of the system in question. We begin with the macroscopic approach.

#### A. Thermomechanical Considerations.

We start with the definition of Eq. (1). For  $V$  and  $T$  independent variables we have at constant volume

$$\Gamma = V \left( \frac{\partial P}{\partial E} \right)_V = V \left( \frac{\partial P}{\partial T} \right)_V / \left( \frac{\partial E}{\partial T} \right)_V . \quad (11)$$

It is a simple problem in calculus to show<sup>4</sup>, by alternatively considering  $V$  as a function of  $P$  and  $T$ , that

$$\left( \frac{\partial P}{\partial T} \right)_V = - \left( \frac{\partial P}{\partial V} \right)_T \left( \frac{\partial V}{\partial T} \right)_P$$

with the result

$$\Gamma = -V \left( \frac{\partial P}{\partial V} \right)_T \left( \frac{\partial V}{\partial T} \right)_P / \left( \frac{\partial E}{\partial T} \right)_V . \quad (12)$$

The derivatives of Eq. (12) each relate to easily performed quasi-static (i.e., non shock) laboratory experiments:

$$\alpha \equiv \text{Thermal expansion coefficient} = \frac{1}{V} \left( \frac{\partial V}{\partial T} \right)_P , \quad (13a)$$

$$\beta \equiv \text{Bulk modulus} = -V \left( \frac{\partial P}{\partial V} \right)_T , \quad (13b)$$

$$C_V \equiv \text{Heat capacity at constant volume} = \left( \frac{\partial E}{\partial T} \right)_V \quad (13c)$$

In terms of  $\alpha$ ,  $\beta$ , and  $C_V$ , Eq. (12) becomes

$$\Gamma = \frac{\alpha \beta}{C_V} V . \quad (14)$$

In the absence of rate dependent effects for the energy deposition problem (e.g., see the discussion relating to Fig. 1) the quasi-static result of Eq. (14) theoretically should give the same  $\Gamma$  as that of the dynamic deposition experiment. In practice such is not always true however, as the energy deposition experiment measures  $\Gamma$  by measuring the pressure of the propagating shock caused by that deposition, and effects such as attenuation and dispersion can easily mask the correct induced pressure (and  $\Gamma$ ).

There is another macroscopic approach of interest, especially since it is related to work done in the next section. It happens that, to a good approximation<sup>5,6</sup>, the thermal properties of solid can be represented by its "characteristic temperature",  $\Theta$ . In that approximation the internal energy can be shown<sup>5,6</sup> to be of the form

$$E = T f\left(\frac{\Theta}{T}\right), \quad (15a)$$

where  $f$  is a function with common form for all solids. Using the first law of thermodynamics, and the definition of the free energy,  $F$ , of a system

$$F = E - TS, \quad (16)$$

where  $S$  is the entropy of the system, one can easily show

$$F = E + T \left( \frac{\partial F}{\partial T} \right)_V. \quad (17)$$

Eq. (17) says that  $F$  and  $E$  have the same functional form

$$F = T g\left(\frac{\Theta}{T}\right), \quad (15b)$$

where again  $g$  is a function with common form for all solids.

Again using the first law of thermodynamics, along with Eq. (16), leads to

$$P = - \left( \frac{\partial F}{\partial V} \right)_T, \quad (18)$$

If we break up the internal energy into a zero temperature contribution, plus a temperature dependent part, then

$$E(V, T) = E_0(V) + E_1(V, T), \quad (19a)$$

$$F(V, T) = E_0(V) + E_1(V, T) - TS(V, T) \equiv F_0(V) + F_1(V, T), \quad (19b)$$

$$F_0(V) = E_0(V), \quad (19c)$$

where Eq. (19c) follows from (19b) and the third law of thermodynamics ( $S \rightarrow 0$  as  $T \rightarrow 0$ ).

$$\therefore P = - \frac{\partial E_0}{\partial V} - \left( \frac{\partial F_1}{\partial V} \right)_T = - \frac{\partial E_0}{\partial V} - \left( \frac{\partial F_1}{\partial \Theta} \right)_T \left( \frac{\partial \Theta}{\partial V} \right), \quad (20)$$

where the second part of Eq. (20) follows from the characteristic temperature approximation as given by Eq. (15b).

Operating upon Eq. (15b) it is a simple matter to show that

$$\left( \frac{\partial F}{\partial \Theta} \right)_T = \frac{\partial g(\Theta/T)}{\partial (\Theta/T)}, \quad (21a)$$

and thus

$$\left. \frac{\partial (F/T)}{\partial (1/T)} \right|_V = \Theta \left( \frac{\partial F}{\partial \Theta} \right)_T. \quad (21b)$$

And using

$$\frac{1}{T} \frac{\partial F}{\partial (1/T)} = \frac{\partial (F/T)}{\partial (1/T)} - F \quad (22a)$$

in Eq. (17) gives

$$\left( \frac{\partial F}{\partial \Theta} \right)_T = \frac{E}{\Theta}. \quad (22b)$$

Substituting Eq. (22b) into Eq. (20) yields

$$P = - \frac{\partial E_c}{\partial V} - \left( \frac{V}{\Theta} \frac{\partial \Theta}{\partial V} \right) \frac{E_1}{V}, \quad (23)$$

where we have applied Eq. (22b) to the temperature dependent contribution alone (subscript 1). For characteristic temperature models such as those of Debye and Einstein, which have a temperature independent  $\Theta$ , comparing Eqs. (23) and (1) gives

$$\Gamma = - \frac{V}{\Theta} \frac{\partial \Theta}{\partial V}. \quad (24)$$

When the Debye characteristic  $\Theta$  is used, Eq. (23) becomes the Debye EOS.

By using the functional details of a given  $\Theta$  model, and experimental Hugoniot data in the form of Eq. (10a), we can thus use Eq. (24) to arrive at numerical values for  $\Gamma$ . Since the experimental Hugoniot data need not be from energy deposition experiments (it could be the result of explosive induced shock measurements), Eq. (24) turns out to be a very useful result. In the next section we shall see that

$$\Gamma = \frac{B_H}{A_H} \quad (25)$$

in terms of the coefficients of Eq. (10a).

## B. Statistical Mechanics Considerations.

We immediately consider a system made up of two weakly interacting subsystems.

As such the formalism is applicable to such systems as fluid filled porous solids<sup>7</sup>, mechanical mixtures (e.g., KDNBF with diatomaceous earth - an explosive mixture sometimes used in electro-explosives devices), and for at least conceptual purposes to electrons and phonons in solids.

Let  $Z$  be the partition function<sup>8</sup> for the system.

$$Z = \sum_{r,s} e^{-\beta E_{rs}} = \sum_{r,s} e^{-\beta(E_r + E_s + \epsilon_{rs})}, \quad (26)$$

where  $E_r$  is a total energy state of one of the subsystems, and  $\epsilon_{rs}$  is the interaction energy between the two subsystems when one subsystem has energy  $E_r$  and the other has energy  $E_s$ . The sum is over all possible energy states consistent with the total energy of the system being constant. The interaction is assumed weak in that  $E_r$  and  $E_s$  are taken to be the subsystem energies in the absence of any interaction. The mean energy and entropy of the system are given in terms of  $Z$  by

$$\bar{E} = - \frac{\partial \ln Z}{\partial \beta}, \quad (27a)$$

$$S = k(\ln Z + \beta \bar{E}). \quad (27b)$$

Eqs. (26) and (27) are sufficient to find the free energy, and thus the pressure, and finally the Grüneisen parameter.

Using  $\beta = (kT)^{-1}$ , and further defining  $\beta \epsilon_{rs} \ll 1$  for a weak interaction, Eq. (26) becomes to first order

$$Z = \sum_{r,s} e^{-\beta E_r} e^{-\beta E_s} (1 - \beta \epsilon_{rs}), \quad (28a)$$

or

$$Z = Z_A Z_B - \beta \sum_{r,s} \epsilon_{rs} e^{-\beta E_r} e^{-\beta E_s}, \quad (28b)$$

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or

$$Z = Z_A Z_B - \beta \sum_{r,s} \epsilon_{rs} e^{-\beta E_r} e^{-\beta E_s}, \quad (28b)$$



where the subsystems are now labeled A and B, and  $Z_A$  is the partition function of subsystem "A" when that subsystem is considered to be an isolated system.

Rewriting Eq. (28b)

$$Z = Z_A Z_B \left[ 1 - \frac{\beta}{Z_A Z_B} \sum_{r,s} \epsilon_{rs} e^{-\beta E_r} e^{-\beta E_s} \right] \equiv Z_A Z_B Z_{AB}. \quad (29a)$$

Thus

$$\ln Z = \ln Z_A + \ln Z_B + \ln Z_{AB}, \quad (29b)$$

and

$$E \equiv \bar{E} = \bar{E}_A + \bar{E}_B - \frac{\partial \ln Z_{AB}}{\partial \beta}. \quad (30)$$

Performing the last indicated operation in Eq. (30) gives

$$-\frac{\partial \ln Z_{AB}}{\partial \beta} = \frac{1}{Z_A Z_B} \sum_{r,s} \epsilon_{rs} e^{-\beta E_r} e^{-\beta E_s} + \frac{\beta}{Z_A Z_B} \sum_{r,s} \epsilon_{rs} \{ \bar{E}_A + \bar{E}_B - E_r - E_s \} e^{-\beta E_r} e^{-\beta E_s}, \quad (31)$$

Similarly applying Eq. (27b) to Eq. (29b) yields

$$S = S_A + S_B + \frac{k\beta^2}{Z_A Z_B} \sum_{r,s} \epsilon_{rs} \{ \bar{E}_A + \bar{E}_B - E_r - E_s \} e^{-\beta E_r} e^{-\beta E_s}, \quad (32)$$

and the free energy takes on the simple form

$$F = F_A + F_B + \frac{1}{Z_A Z_B} \sum_{r,s} \epsilon_{rs} e^{-\beta E_r} e^{-\beta E_s}, \quad (33a)$$

or

$$F = F_A + F_B + F_{AB}, \quad (33b)$$

which defines  $F_{AB}$ . Although we could have guessed the result Eq. (33b) for two weakly interacting subsystems, Eq. (33a) allows one to calculate  $F_{AB}$  when one has a model for the subsystems and their interaction.

### C. General Applications.

Consider the case (e.g., a mixture of weakly interacting gases) when the subsystems can be considered as occupying the same volume. Then

$$P = - \left( \frac{\partial F}{\partial V} \right)_T = P_A + P_B - \left( \frac{\partial F_{AB}}{\partial V} \right)_T, \quad (34a)$$

or

$$PV = P_A V + P_B V - V \left( \frac{\partial F_{AB}}{\partial V} \right)_T, \quad (34b)$$

and

$$\Gamma E = \Gamma_A \bar{E}_A + \Gamma_B \bar{E}_B - V \left( \frac{\partial F_{AB}}{\partial V} \right)_T, \quad (34c)$$

where Eq. (1) has been used. The result Eq. (34c) is quite significant in that, even for non-interacting subsystems, it says that the Gruneisen parameter is a weighted average of the Gruneisen parameters of the subsystems. Eq. (34c) can, for example, be applied to problems such as the absorption of a high intensity laser beam by the air-vapor mixture in front of an irradiated solid surface<sup>9</sup>.

Let us now apply Eq. (33b) to the case of a mechanical mixture, the subsystems (components) occupy volume separately. For brevity we neglect the interaction energy between the subsystems. Thus

$$F = F_A + F_B, \quad (35a)$$

$$V = V_A + V_B. \quad (35b)$$

$$\therefore P = - \left( \frac{\partial F}{\partial V} \right)_T = - \left( \frac{\partial F_A}{\partial V_A} \right)_T \frac{dV_A}{dV} - \left( \frac{\partial F_B}{\partial V_B} \right)_T \frac{dV_B}{dV}. \quad (36a)$$

Eq. (36a) contains the realistic assumption that  $F_A$  does not depend upon  $V_B$ , and similarly for  $F_B$ .

$$P = P_A \frac{dV_A}{dV} + P_B \frac{dV_B}{dV}, \quad (36b)$$

$$PV = \Gamma E = P_A V \frac{dV_A}{dV} + P_B V \frac{dV_B}{dV}, \quad (36c)$$

and then

$$\Gamma = \Gamma_A \left( \frac{\bar{E}_A}{E} \right) \left( \frac{V}{V_A} \frac{dV_A}{dV} \right) + \Gamma_B \left( \frac{\bar{E}_B}{E} \right) \left( \frac{V}{V_B} \frac{dV_B}{dV} \right). \quad (36d)$$

The result for a mixture of two non-interacting monatomic gases can be re-written, from Eq. (34c) and (2), as

$$\Gamma = \Gamma_A \left( \frac{N_A}{N} \right) + \Gamma_B \left( \frac{N_B}{N} \right) \quad (37)$$

an expected result. Eq. (36d) is thus reasonably explained as Eq. (37) plus a volume effect. We note that, for  $N = N_A + N_B$ ,  $\Gamma_A \leq \Gamma_B$  assumed, we have  $\Gamma$  varying linearly between the extremes  $\Gamma_A$ ,  $\Gamma_B$ .

For the special case of components A and B undergoing equal strain (containing zero strain in both components as a special case), the solid mixture behaves as the gaseous mixture of Eq. (37). For equal strains

$$\frac{dV_A}{V_A} = \frac{dV_B}{V_B}, \quad (38a)$$

but

$$\frac{dV}{V} = \frac{V_A}{V} \left( \frac{dV_A}{V_A} \right) + \frac{V_B}{V} \left( \frac{dV_B}{V_B} \right). \quad (38b)$$

$$\therefore \frac{dV}{V} = \frac{dV_A}{V_A} = \frac{dV_B}{V_B}. \quad (38c)$$

Using Eq. (38c) in Eq. (36d) yields

$$\Gamma = \Gamma_A \left( \frac{\bar{E}_A}{E} \right) + \Gamma_B \left( \frac{\bar{E}_B}{E} \right), \quad (39)$$

which has  $\Gamma$  varying linearly between  $\Gamma_A$  and  $\Gamma_B$  since  $E = \bar{E}_A + \bar{E}_B$ . We note that if the strains are equal the constant volume definition of Eq. (1) holds even for energy deposition pulse duration large compared to the time for acoustic propagation across a component particle; if the strains are equal they will be zero until a pressure wave due to inhomogeneous deposition reaches the local observation point. Thus for equal strains, which contains a statement concerning elastic constants, Eq. (39) appears to hold for energy deposition times both large and small compared to acoustic propagation times.

There is another special case which can be treated analytically. If the component strains are not equal (thus allowing local pressure relief), then for energy deposition times large compared to the time for acoustic propagation across a particle diameter

$$P = P_A = P_B, \quad (40a)$$

and 
$$PV = P_A V = P_B V_A \left( \frac{V}{V_A} \right). \quad (40b)$$

$$\therefore \Gamma^* = \Gamma_A^* \left( \frac{\bar{E}_A}{E} \right) \left( 1 + \frac{V_B}{V_A} \right), \quad (40c)$$

where  $\Gamma_A^*$  indicates the "effective" Grüneisen constant<sup>10</sup> since the constant volume conditions of Eq. (1) are not satisfied. Alternatively, Eq.(40a) can be used to substitute the compressibility,  $K$ , for the volume terms in Eq. (36d).

$$K \equiv -\frac{1}{V} \frac{dV}{dP}. \quad (41)$$

When Eqs. (40a) and (41) are used with Eq. (36d) one gets

$$\Gamma^* = \Gamma_A^* \frac{\kappa_A \bar{E}_A}{\kappa E} + \Gamma_B^* \frac{\kappa_B \bar{E}_B}{\kappa E} . \quad (42)$$

If the approximations of Eqs. (38c) or (40a) do not hold, then Eq. (36d) must be used along with a model for  $V_A = V_A(V)$  and  $V_B = V_B(V)$ .

#### IV. SOLIDS

In this section we consider the individual contributions to the Gruneisen parameter of a solid. We consider only metals because

- a. The results for metals contain the results for dielectrics as a special case.
- b. As mentioned in the introduction we have a strong interest in electromechanical effects in metals.
- c. For the most part, aerospace materials are composed of metals, dielectrics, or their mixture (e.g., composites). Further, in present technology, aerospace applications stand the best chance of seeing an environment where a knowledge of the Gruneisen parameter is important.

We break up the problem into electron and phonon\* contributions while neglecting the volume dependent zero degree Kelvin cohesive energy in keeping with the constant volume definition of Eq. (1). Thus

$$F(V, T) = F_e(V, T) + F_p(V, T) + \mathcal{E}_{ep}(V, T), \quad (43)$$

where the subscript e denotes electrons, p denotes phonons, and  $\mathcal{E}_{ep}(V, T)$  the contribution to the temperature dependent part of the free energy from electron-phonon interactions. Even though the interaction term appears to be crucial to the electromechanical problem, for the Gruneisen parameter problem one gets good agreement with experiment by considering the electrons and phonons to comprise non-interacting gases which occupy the same volume. Thus we neglect the interaction term and note that Eq. (34c) applies.

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\*A phonon is a quantized unit of lattice vibration in a normal mode system where the vibrations behave as a set of non-interacting oscillators.

### A. Electrons.

From Zharkov and Kalinin<sup>3</sup> we have for the electron system

$$E_e(V, T) = \frac{2\pi^3}{3h^2} m \left( \frac{3N}{\pi} \right)^{1/3} V^{2/3} (kT)^2, \quad (44a)$$

$$S_e(V, T)T = \frac{4\pi^3}{3h^2} m \left( \frac{3N}{\pi} \right)^{1/3} V^{2/3} (kT)^2, \quad (44b)$$

$$\therefore F_e(V, T) = - \frac{2\pi^3}{3h^2} m \left( \frac{3N}{\pi} \right)^{1/3} V^{2/3} (kT)^2, \quad (44c)$$

$$\therefore P_e(V, T)V = \frac{4\pi^3}{9h^2} m \left( \frac{3N}{\pi} \right)^{1/3} V^{2/3} (kT)^2, \quad (44d)$$

$$\therefore \Gamma_e = \frac{2}{3} \quad (44e)$$

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In the above  $m$  is the mass of an electron,  $N$  is the number of electrons in the system, and  $h$  is Planck's constant ( $h = 6.63 \times 10^{-27}$  erg seconds). Comparing Eq. (44e) with Eq. (3) shows that the electrons behave as if they were an ideal gas. Eq. (44e) can also be derived on the more general grounds<sup>11</sup> of thermodynamics with a Planck distribution, which then gives Eq. (3) as a special case.

For the phonon contribution we break the problem up into low and high temperature regimes. The reason for the separation is the radically different functional dependence of the energy and entropy in the two regimes. The formulae for both regimes is from reference eleven.

### B. Phonons-Low Temperatures.

$$E_p(V, T) = V \frac{\pi^2 (kT)^4}{10 (\hbar \bar{c})^3}, \quad (45a)$$

$$S_p(V, T)T = V \frac{2\pi^2 (kT)^4}{15 (\hbar \bar{c})^3}, \quad (45b)$$

$$\therefore F_p(V, T) = -V \frac{\pi^2 (kT)^4}{30 (\hbar \bar{c})^3}, \quad (45c)$$

$$VP_p(V, T) = - V^2 \frac{\pi^2 \hbar (kT)^4}{10(\hbar \bar{c})^4} \left( \frac{\partial \bar{c}}{\partial V} \right)_T \quad (45d)$$

$$\therefore P_p = - \frac{V}{\bar{c}} \left( \frac{\partial \bar{c}}{\partial V} \right)_T. \quad (45e)$$

In arriving at Eq. (45d) we have treated the factor of  $V$  in Eq. (45c) as a constant. This is because the factor of  $V$  tells how much material is present without saying anything about the structure, while the pressure is determined by the structure. The electron subsystem was treated in the same way in that the number of electrons,  $N$ , was taken to be constant.  $\hbar \equiv 2\pi \hbar$ , and  $\bar{c}$  is defined by

$$\frac{1}{\bar{c}^3} = \frac{1}{c_l^3} + \frac{2}{c_t^3}, \quad \text{Reproduced from best available copy.} \quad (46)$$

where  $c_l$  is the longitudinal sound velocity, and  $c_t$  is the transverse sound velocity.

### C. Phonons - High Temperature.

At high temperatures the sums over the vibrational states become greatly simplified <sup>11</sup> with the result

$$\bar{E}_p(V, T) = 3NkT, \quad (47a)$$

$$S_p(V, T)T = 3NkT \left[ \ln kT - \ln \hbar \bar{\omega} + 1 \right], \quad (47b)$$

where  $N$  is the number of atoms (for brevity we consider a monatomic lattice) and  $\bar{\omega}$  is some mean vibrational frequency associated with the lattice structure. It is obvious from Eqs. (47) that only the entropy contributes to the pressure.

$$\therefore PV = - 3NkTV \left( \frac{\partial \bar{\omega}}{\partial V} \right)_T, \quad (48a)$$

$$\therefore P_p = - \frac{V}{\bar{\omega}} \left( \frac{\partial \bar{\omega}}{\partial V} \right)_T. \quad (48b)$$



If we now use the relationship from lattice dynamics <sup>4</sup>

$$\bar{\omega} = q \bar{c}, \quad (49)$$

where  $q$  is the magnitude of the state vector  $\vec{q}$ , then we see that the high and low temperature regimes yield the same result

$$\Gamma_p = - \frac{v}{c} \left( \frac{\partial \bar{c}}{\partial v} \right)_T. \quad (45e)$$

A few words are in order regarding the difference between  $\Gamma_e = 2/3$  and the result Eq. (45e). The energy of an individual phonon can be written

$$\epsilon = p \bar{c}, \quad (50)$$

where  $p$  is "momentum" associated with the phonon state. The form of Eq. (50) is radically different from that for the energy of a gas molecule in an ideal gas, namely,

$$\epsilon = \frac{p^2}{2m} \quad (51)$$

where  $m$  is the molecular mass. Thus the difference between  $\Gamma_e$  and  $\Gamma_p$  is explained by the non-gas-like behavior of the phonon system.

From the discussion of section II of this report we can assume that  $\Gamma_p$  is of order unity (we shall actually determine it later in this section), and apply Eq. (34c) with the interaction term taken as zero to judge the relative importance of the electron and phonon contributions. We need only consider the ratio  $E_e(V, T)/E_p(V, T)$ . The ratio is calculated below for  $T = 300^\circ\text{K}$  (high temperature limit), for the number of electrons equal to the number of vibrating atoms (i.e., a metal), and an electron density of  $10^{21}$  per  $\text{cm}^3$ . The low temperature ratio is found to be unity at  $10^\circ\text{K}$ .

$$\text{High Temp.:} \quad \frac{E_e(v, T)}{E_p(v, T)} = \frac{2\pi^2 m k T}{3h^2 \left(\frac{3n}{\pi}\right)^{1/3}} \sim \frac{1}{2c}. \quad (52)$$

$$\text{Low Temp.:} \quad \frac{E_e(v, T)}{E_p(v, T)} = \frac{5mh \left(\frac{3n}{\pi}\right)^{1/3} \bar{c}^3}{6\pi^2 (kT)^2} \sim 1. \quad (53)$$

In Eqs. (52) and (53) we have taken a typical solid acoustic velocity,  $\bar{c}$ , of  $5 \times 10^5$  cm per sec, and  $n$  is the density of electrons in the conduction band.

We thus see that at room temperature only the phonon contribution is of importance, while below  $10^\circ\text{K}$  and above  $6,000^\circ\text{K}$  the electron and phonon contributions are of equal magnitude.

Let us now calculate  $\Gamma_p$  from experimental shock wave data. To do so we use Eqs. (10b), (45e), and

$$\bar{c}(p) \equiv \sqrt{\frac{\partial P_H}{\partial \rho}}. \quad (54)$$

In using Eq. (10b) for calculating  $\Gamma_p$  we are assuming that the density variation of  $P_H$  is primarily due to phonon effects. Since for metals the electron "sound" (i.e., Fermi) velocity<sup>5</sup> is of order  $10^7$  cm per sec, which is about a factor of 20 larger than the observed propagation velocity<sup>3</sup> for a weak shock, the assumption would appear to be valid. The assumption is also consistent with the discussion accompanying Eq. (52). The arithmetic is straightforward:

$$\bar{c}(p) = \left(\frac{A_H}{P_0}\right)^{1/2} \sqrt{1 + \frac{2B_H}{A_H} \left(\frac{p}{P_0} - 1\right)}, \quad (55a)$$

$$\bar{c}(p) \cong \bar{c}_e \left\{ 1 + \frac{B_H}{A_H} \left(\frac{p}{P_0} - 1\right) \right\}, \quad (55b)$$

where

$$\frac{B_H}{A_H} \left(\frac{p}{P_0} - 1\right) \ll 1 \quad (55c)$$

is assumed, and

$$\bar{c}_o \equiv \sqrt{\frac{A_H}{\rho}} \quad (55d)$$

$$\therefore \frac{V}{\bar{c}} \frac{\partial \bar{c}}{\partial V} \approx \frac{V}{\bar{c}} \frac{\partial}{\partial V} \left\{ \bar{c}_o \frac{B_H}{A_H} \left( \frac{V_o}{V} - 1 \right) \right\} = - \frac{V}{\bar{c}} \left\{ \bar{c}_o \frac{B_H}{A_H} \frac{V_o}{V^2} \right\} \quad (55e)$$

$$\therefore \Gamma_P = \frac{B_H}{A_H} > \quad (25)$$

where we have evaluated Eq. (55e) at  $V = V_o$ .

Thus by using shock wave data generated by non-energy deposition experiments (e.g., flyer experiments), we arrive at a Grüneisen parameter for use in energy deposition induced shock work.

Table I below compares  $\Gamma_P$  for three metals from Eq. (25) with the Grüneisen parameter based upon the thermodynamic quantities of Eq. (14).

TABLE I  
COMPARISON OF SHOCK AND THERMODYNAMIC DATA

Aluminum	2.17 <sup>a</sup>	2.09 <sup>b</sup>
Copper	2.00	1.98
Lead	2.78	2.46

<sup>a</sup>. Hugoniot data from Table IV of reference 2.

<sup>b</sup>. Values from L. V. Al'tshuler et al, Soviet Phys. J. E. T. P. 11, 573 (1960)

The agreement shown in Table I is rather good, and is typical of the agreement which one would get for a more extensive list of metals.

#### D. Volume Dependence.

It is also possible to include a volume dependence for  $\Gamma$  in the calculation leading to Eq. (25). One need simply not evaluate Eq. (55e) at  $V = V_0$ , and for higher volume terms simply add on additional higher order terms to Eq. (10b). The volume dependence of  $\Gamma$  is indeed quite important as can be seen from the data for aluminum (from ref. 2) given below

$$\Gamma(V) = 2.13 - 7.2 \left( \frac{\rho}{\rho_0} - 1 \right) + 25 \left( \frac{\rho}{\rho_0} - 1 \right)^2. \quad (56)$$

The volume terms of Eq. (56) contribute twenty percent to  $\Gamma(V)$  for  $\rho/\rho_0 = 1.2$ , and the result would even be more impressive if the second and third terms on the right of Eq. (56) were not of opposite sign.

We should at this time comment on the physics of such a volume dependence, namely where it comes from. In Eq. (8) let  $b$  have a volume dependence (to be explained shortly), and expand  $\Gamma$  as a function of volume for a small volume change about an initial volume  $V_0$ . Thus

$$\Gamma(V_0 + \Delta V) = \Gamma(V_0) + \frac{2}{3} \left[ \left. \frac{db}{dV} \right|_{V_0} \frac{V_0}{b(V_0)} - 1 \right] \frac{\Delta V}{V_0}. \quad (57)$$

There are two contributions to Eq. (57). The first involves the volume change alone (i.e.,  $\Delta V/V_0$ ) and is explained by the observation that as the volume increases the molecular volume becomes less important; since pressure is a measure of the statistics of particle interactions with the walls of an imaginary container, the presence of other particles which occupy volume effects those statistics.

The second contribution to Eq. (57) is proportional to  $(db/dV)$ , and says that the molecular volume is dependent upon the density of neighboring molecules. This is understood by saying that where the volume dependence of the Grüneisen parameter is concerned, an exponential repulsive interaction is more appropriate than a hard sphere interaction. Indeed, by using such an exponential repulsive term Sirdeshmukh and Rao<sup>12</sup> were able to derive  $\Gamma$  (and its volume dependence) for some crystals of the fluorite structure.

Before leaving this section we wish to mention another method for arriving at the volume dependence of the Grüneisen parameter, namely the Dugdale - MacDonald formula<sup>2,3</sup>. That formula is

$$\Gamma = - \frac{V}{2} \frac{\partial^2 (PV^{2/3}) / \partial^2 V}{\partial (PV^{2/3}) / \partial V} - \frac{1}{3}. \quad (58)$$

Eq. (57) is derived for an isotropic solid at 0°K by essentially taking the third derivative with respect to volume of the lattice potential energy (which is assumed to be only a function of the atomic coordinates). Eq. (58) works because, if  $\phi$  is the potential energy, and if the potential is quadratic in the lattice parameter, then

$$\frac{\partial^3 \phi}{\partial V^3} \propto \frac{\partial \omega}{\partial V}, \quad (59)$$

$(\omega/2\pi)$  being the vibrational frequency of the 0°K unoccupied vibrational states.

The success of Eq. (58) is understood as the right hand side of the proportionality, Eq. (59), is the prime ingredient of Eq. (48b).

Equation (58) is one of a whole class of relations<sup>3</sup> which utilize macroscopic PV data to arrive at a Grüneisen constant. All such relations, including Eq. (25)

and the procedure of this section used in deriving it, assume an isotropic solid. This means that the Hugoniot data used in arriving at  $\Gamma$  must be in a range where pressure is large compared to the yield stress (stress for the onset of plasticity). This is because most of the Hugoniot data is arrived at via one-dimensional strain experiments (e.g., flyer plate with planar geometry), with the result that solids are not in an isotropic state in the region of the yield stress<sup>13</sup>. The yield stress of a typical metal lies between a few kilobars and tens of kilobars.

#### E. Ultrasonic Data.

It frequently happens that Hugoniot data does not exist for a material. In such a case ultrasonic data can often be used to evaluate  $\Gamma$ . What is needed experimentally is the usual apparatus for measuring elastic constants, and additional equipment which allows for those measurements to be carried out as a function of pressure (a maximum pressure of a few kilobars is usually sufficient). The theory is relatively simple. If  $\sigma_{ij}$  is the stress tensor, then

$$\sigma_{ij} = C_{ijkl} \epsilon_{kl} + D_{ijklmn} \epsilon_{kl} \epsilon_{mn}, \quad (60)$$

where  $\epsilon_{kl}$  is the strain tensor, and C and D are the second order and third order elastic constants respectively<sup>5</sup>. For a cubic crystal undergoing an isotropic volume change,  $P = \sigma_{11} = \sigma_{22} = \sigma_{33}$ , and

$$\epsilon_{11} + \epsilon_{22} + \epsilon_{33} = \frac{V}{V_0} - 1. \quad (61)$$

Thus the constants C and D can be related to the constants  $A_H$  and  $B_H$  of Eq. (10a).

The D's are obtained through the pressure variation of the sound velocity. Eqs.

(60) and (61) when combined are of the same form as Eq. (10b).

While the Hugoniot EOS approach to the Grüneisen constant is valid for isotropic materials (e.g., polycrystalline metals) it is not valid for anisotropic single crystals<sup>14</sup>. This is because the Grüneisen parameter is really a second rank tensor and should be defined by

$$\Gamma_{ij} = V \left( \frac{\partial \sigma_{ij}}{\partial E} \right)_V. \quad (62)$$

Eq. (1) is then seen to be Eq. (62) applied to isotropic solids. Since the C's and D's of Eq. (60) give all of the necessary information for evaluating the various elements of  $\Gamma$ , the ultrasonic method is the superior tool. Indeed application of the hydrostatic Hugoniot method to an anisotropic crystal such as Zirconium (hexagonal close packed)<sup>15</sup> is now known to give  $\Gamma$  values inconsistent with thermomechanical results because of anisotropic strain effects.

We are essentially making two points here. First, that in addition to the usual Hugoniot data, shock wave physicists have relevant ultrasonic data at their disposal. Second, that for anisotropic media, only the ultrasonic data is applicable.

## V. DISCUSSION AND SUMMARY

The main intent of this report has been to give the shock wave physicist, the prime applied user of the Grüneisen parameter, an understanding of the physics involved in that parameter as well as the techniques in calculating it from related data.

In this closing section we make a few points which we hope will serve to tie up any loose strings. We also mention some relatively esoteric physics contributions to which at this time although interesting in their own right are relatively unimportant in applications.

Within the main text of this report we failed to mention the connection between the characteristic  $\Theta$  Grüneisen parameter

$$\Gamma = - \frac{V}{\Theta} \frac{\partial \Theta}{\partial V} , \quad (24)$$

and that arrived at through phonon considerations

$$\Gamma_P = - \frac{V}{\bar{\omega}} \frac{\partial \bar{\omega}}{\partial V} . \quad (48b)$$

The connection is that the characteristic temperature  $\Theta$  allows a characteristic frequency,  $\omega_m/2\pi$ , to be defined by

$$\hbar \omega_m \equiv k \Theta . \quad (63)$$

Thus if  $\omega_m$  has the same volume dependence as  $\bar{\omega}$  the two approaches will give the same result.

At high temperatures the detailed definition of  $\bar{\omega}$  for a monatomic lattice is given by<sup>11</sup>

$$\log \bar{\omega} \equiv \frac{1}{3N} \sum_s \log \omega_s , \quad (64)$$



where  $\alpha$  is an index over the possible vibrational states of the solid. The allowed number of  $\omega_\alpha$  increase as  $\omega_\alpha^3$  (a relationship known as the density of states<sup>11</sup>) so that we expect  $\bar{\omega} \approx \omega_m$  as  $\omega_m$  is associated with the highest possible energy state for a characteristic  $\Theta$  model. Thus we are led to expect equality in the two approaches at high temperatures.

At low temperatures a probability factor is present (relating to the probability that a given state is occupied<sup>11</sup>, in the equivalent of Eq. (63). This probability factor serves to reduce the ratio  $\bar{\omega}/\omega_m$  to zero as the temperature goes to zero. Since  $\bar{\omega}$  is no longer close to  $\omega_m$  they do not necessarily have the same volume dependence - and indeed for many materials they do not have the same volume dependence.

Germanium offers a good example of the complications introduced by low temperatures. Experimentally the thermomechanical Gruneisen parameter, Eq. (14), of Germanium is negative<sup>14</sup> in the region between 10°K and 40°K. Yet  $\frac{\partial \omega_m}{\partial V}$  is determined only by the crystal structure and remains positive in that same region. That low temperatures present difficulties should not be surprising as we have already seen that electron contribution to  $\Gamma$  is equal to the phonon contribution at about 10°K. Those interested in further details of the electronic contributions to are referred to the work of Wallace<sup>16</sup>.

There is another conceptual approach to the physics of the Gruneisen parameter which is very much worth mentioning. It is correct to think of phonons and electrons as representing thermally excited states of a crystal lattice in the same sense as kinetic energy is the thermal excitation of an ideal gas. But there are a host of other possible excited states in a crystal lattice which are more or less important

depending upon temperature, volume, and the types of atoms making up the crystal in question. Some of these thermal excitations are listed below.

- MAGNONS<sup>5</sup> - elementary propagating magnetic waves in paramagnetic materials.
- EXCITONS<sup>5</sup> - coupled electron-hole pair states.
- ROTONS<sup>17</sup> - elementary rotational states in a superfluid.
- ELECTRON-ELECTRON PAIRS<sup>5</sup> - the elementary excitations responsible for superconductivity.
- HELICONS<sup>18</sup> - an electromagnetic excitation in a solid state plasma.

The point is that each type of thermal excitation represents a mechanism for storing energy. Like electrons and phonons the energy dependence of the partial pressures contributed by those more exotic excitations must be considered in evaluating the Gruneisen parameter.

The effect of the more exotic excitations is presently not a serious question for the applied shock wave physicist. However, as environments change and previously strange materials are introduced into technology .....

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34 - 45

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46

47

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49

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